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Unified origin of axion and monopole dark matter, and solution to the domain-wall problem

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We propose a scenario in which the spontaneous breakdown of the Peccei-Quinn symmetry leads to monopole production. Both the axion and the monopole contribute to dark matter, and the Witten effect on the axion mass is a built-in feature. In the Kim-Shifman-Vainshtein-Zakharov-type axion model, seemingly different vacua are actually connected by the hidden gauge symmetry, which makes the axionic string unstable and separate into two Alice strings. The Alice string is attached to a single domain wall due to the QCD instanton effect, solving the domain-wall problem. This is in the same spirit of the Lazarides-Shafi mechanism, although the discrete Peccei-Quinn symmetry is not embedded into the center of the original gauge symmetry. In the Dine-Fischler-Srednicki-Zhitnitsky-type axion model, the domain-wall problem is avoided by the Witten effect. If the Peccei-Quinn symmetry is explicitly broken by a small amount, monopoles acquire a tiny electric charge and become minicharged dark matter. Interestingly, the quality of the Peccei-Quinn symmetry is closely tied to darkness of dark matter.

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I. INTRODUCTION

The origin of dark matter (DM) remains a mystery in both cosmology and particle physics. Its longevity may be due to conserved charge, light mass, or very weak interactions to the standard model. One example is a hidden monopole of which the stability is guaranteed by its topological charge [1–3]. Another is the QCD axion [4–7] or axionlike particles of which the mass and couplings are suppressed by the decay constant.

Monopoles arise at a spontaneously symmetry breaking (SSB) of non-Abelian symmetry, if the vacuum manifold has a nontrivial second homotopy group. For instance, an $\mathrm{SU}(2)_H$ gauge symmetry is spontaneously broken down to $\mathrm{U}(1)_H$ in the case of a simple 't Hooft-Polyakov monopole [8,9], and its cosmological implication was studied in Refs. [2,3].

The monopole abundance is determined by the correlation length at the SSB, which depends on the detailed dynamics of the phase transition. If the phase transition is of the first order, monopoles can be produced when the expanding bubbles collide. In the case with a Coleman-Weinberg potential, a monopole with mass of $\mathcal{O}(10^{10})$ GeV can explain the observed DM density for a hidden gauge coupling $g_H = \mathcal{O}(0.1)$ [3]. On the other hand, if the phase transition is of the second order, the monopole mass should be of $\mathcal{O}(10^2)$ TeV since its abundance is much larger [1–3].

The QCD axion is a Nambu-Goldstone boson associated with the spontaneous breakdown of a global U(1) Peccei-Quinn (PQ) symmetry. The PQ breaking scale sets the axion decay constant f_a, and they are comparable to each other in a simple setup. The QCD axion and axionlike particles have been extensively studied in the literature (see, e.g., Refs. [10–13] for recent reviews). The observed neutrino burst duration of SN1987A implies that the axion decay constant cannot be smaller than 4×10^8 GeV [14]. In the

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early Universe, the axion is produced by the misalignment mechanism [15–17]. If the PQ symmetry is already broken during inflation, this leads to an upper bound, $f_a \lesssim 10^{12}$ GeV, barring fine-tuning of the initial displacement.[1] If the PQ symmetry gets spontaneously broken after inflation, on the other hand, axions are produced by the decay of domain walls and cosmic strings. In this case, the axions explain the observed DM abundance if the PQ breaking scale is of $\mathcal{O}(10^{10})$ GeV [20] or slightly larger [21].

Recently, it was pointed out that there is an interesting interplay between the QCD axion and hidden monopoles [22–24]; if the QCD axion is coupled to a hidden $\mathrm{U}(1)_H$ under which the hidden monopole is charged, the Witten effect induces an effective axion mass in the presence of monopoles [25,26]. The effective axion mass squared is proportional to the monopole number density, which decreases as the Universe expands. Therefore, the Witten effect is important only in the early Universe and does not spoil the PQ mechanism in the present vacuum. Interestingly, one can solve cosmological problems of the QCD axion such as the overabundance/isocurvature problem and the domain-wall problem by making use of the Witten effect [22–24].

The proximity of the two symmetry breaking scales in the monopole and axion DM scenarios implies that they may have a common origin. In this paper, we pursue this possibility and build a simple model that unifies the origin of the QCD axion and monopole. The Witten effect on the axion mass is a built-in feature of this model. We also find that the domain-wall problem of the QCD axion can be avoided either by the Witten effect or by the way similar to the Lazarides-Shafi mechanism [27]. Interestingly, the high quality of the PQ symmetry is closely related to the darkness of DM. In other words, if one introduces small PQ symmetry breaking terms, monopoles acquire a tiny electric charge and become minicharged DM, which can be searched for by, e.g., direct DM search experiments.

The rest of this paper is organized as follows. In the next section, we provide a simple model in which the PQ symmetry breaking generates 't Hooft-Polyakov monopoles in a hidden sector. We study the DM abundances in Sec. IV and observational signatures in Sec. V. The last section is devoted to discussion and conclusions.

II. MODEL

We introduce a hidden $\mathrm{SU}(2)_H$ gauge symmetry and an adjoint complex scalar field Φ. We further impose a global $\mathrm{U}(1)_{\mathrm{PQ}}$ symmetry on Φ, which will play a role of the PQ scalar. We assume that $\mathrm{U}(1)_{\mathrm{PQ}}$ symmetry is anomalous under $\mathrm{SU}(3)_c$ gauge symmetry for the PQ mechanism to work.

The potential for Φ is given by

$$V(\Phi) = -\frac{M^2}{2}\mathrm{Tr}\Phi\Phi^\dagger + \frac{\lambda_1}{4}(\mathrm{Tr}\Phi\Phi^\dagger)^2 + \frac{\lambda_2}{4}|\mathrm{Tr}\Phi^2|^2, \qquad (1)$$

where Tr represents the trace over the $\mathrm{SU}(2)_H$ indices, $\lambda_1 + \lambda_2$ is taken to be positive for the potential to be bounded below, and we assume $M^2 > 0$ and $\lambda_2 < 0$ to obtain the successful spontaneous breaking of the PQ symmetry and the monopole solution.[2] The vacuum expectation value (VEV) of Φ is given as

$$\langle\Phi\rangle = \frac{v}{\sqrt{2}}\begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}, \qquad v = \frac{M}{\sqrt{\lambda_1 + \lambda_2}}. \qquad (2)$$

This spontaneously breaks $\mathrm{SU}(2)_H$ gauge symmetry down to $\mathrm{U}(1)_H$. At the same time, the $\mathrm{U}(1)_{\mathrm{PQ}}$ global symmetry is spontaneously broken, leading to the presence of the QCD axion. The SSB scale v therefore sets both the monopole mass and the PQ breaking scale (i.e., axion decay constant).

In addition to the $\mathrm{U}(1)_H$ symmetry, the above VEV of Φ respects a Z_2 symmetry, which is generated by $\exp(i\pi T_2)\exp(i\pi Q_{\mathrm{PQ}})$, where T_2 is a broken generator of $\mathrm{SU}(2)_H$ and Q_{PQ} is the generator of $\mathrm{U}(1)_{\mathrm{PQ}}$. This Z_2 transformation does not commute with the $\mathrm{U}(1)_H$ gauge transformation. Therefore, the unbroken symmetry in an era between the PQ phase transition and the QCD phase transition is $\mathrm{U}(1)_H \rtimes Z_2$ symmetry.[3] Our model is a variant of the Alice electrodynamics [30,31].

The structure of the vacuum can be parametrized by using a real unit vector $\vec{a}$ in a three-dimensional space and a phase ϕ in $\mathrm{U}(1)_{\mathrm{PQ}}$. However, $(\vec{a}, \phi)$ and $(-\vec{a}, \phi + \pi)$ are identical points because of the Z_2 symmetry described above. Therefore, the global structure of the vacuum is $(S_2 \times S_1)/Z_2$. As we expected, this manifold has a nontrivial second homotopy group. This is the monopole solution associated with the $\mathrm{SU}(2)_H \to \mathrm{U}(1)_H$ breaking. We also have two types of cosmic strings. One is the usual axionic cosmic string in which the phase of Φ is rotated by $\mathrm{U}(1)_{\mathrm{PQ}}$ transformation around the cosmic string. The other is a cosmic string in which the phase of Φ is rotated by π due to $\mathrm{U}(1)_{\mathrm{PQ}}$ transformation and then its sign is changed by the $\mathrm{SU}(2)_H$ gauge transformation. For example, we can consider the configuration around the string,

[1] The upper bound is greatly relaxed for low-scale inflation with the Hubble parameter comparable to or less than the QCD scale [18,19].

[2] For $\lambda_2 > 0$, the $\mathrm{SU}(2)_H$ gauge symmetry is completely broken, and no monopole solution exists. In the marginal case, monopoles are created at the first phase transition, and they survive for some time until the residual $\mathrm{U}(1)_H$ gauge symmetry is spontaneously broken by the second phase transition.

[3] The same symmetry breaking pattern has been studied in different contexts [28,29].

$$\langle\Phi\rangle = \frac{ve^{i\theta/2}}{\sqrt{2}} \times \begin{pmatrix} \cos(\theta/2) & \sin(\theta/2) \\ \sin(\theta/2) & -\cos(\theta/2) \end{pmatrix}, \quad (3)$$

where θ is the azimuthal angle of the cylindrical coordinate around the cosmic string. This is the so-called Alice string [32,33], in which the cosmic string consists of both global symmetry and local symmetry. As we will see later, this Alice string has a crucial role in the evolution of the string-wall system.

The scalar field Φ has 6 degrees of freedom: two degrees are eaten by the longitudinal component of the $W_H^\pm$ bosons, two form a massive $U(1)_H$ charged scalar $\Phi^\pm$, one is the QCD axion a, and the remaining one is a massive Higgs boson ϕ. Here, the superscripts denote the $U(1)_H$ charge. The masses of these heavy gauge scalar bosons are

$$m_{W_H^\pm} = \sqrt{2} g_H v, \quad (4)$$

$$m_\phi^2 = 2(\lambda_1 + \lambda_2) v^2, \quad (5)$$

$$m_{\Phi^\pm}^2 = 2|\lambda_2| v^2, \quad (6)$$

where g_H is the $SU(2)_H$ gauge coupling.

We will also introduce couplings of Φ with the PQ quarks or hidden fermions. After integrating them out [as well as the standard model (SM) quarks in the Dine-Fischler-Srednicki-Zhitnitsky (DFSZ)-type axion model], we obtain the axion coupling to gluons and hidden photons at low energy,

$$N_{\rm DW}^{(c)} \frac{\alpha_s}{8\pi} \frac{a}{v} G_{a\mu\nu} \tilde{G}^{a\mu\nu} + N_{\rm DW}^{(H)} \frac{\alpha_H}{8\pi} \frac{a}{v} F'_{\mu\nu} \tilde{F}'^{\mu\nu}, \quad (7)$$

where $G_{a\mu\nu}$ and $F'_{\mu\nu}$ are the field strengths of gluons and the $U(1)_H$ hidden photon, respectively; their duals are shown with tildes; α_s and $\alpha_H = g_H^2/4\pi$ denote the strong coupling constant and the fine-structure constant of $U(1)_H$, respectively; and $N_{\rm DW}^{(c)}$ and $N_{\rm DW}^{(H)}$ are the so-called domain-wall numbers specified below.

The axion decay constant f_a is defined by $f_a = v/N_{\rm DW}^{(c)}$. The axion coupling to gluons gives rise to the axion potential so that the axion is stabilized at the strong CP-conserving point, solving the strong CP problem. The coupling to hidden photons usually does not give any mass to the axion, but in the presence of hidden monopoles, it generates the effective axion mass that depends on the monopole number density via the Witten effect [25,26].

A. KSVZ axion model

Now, we shall couple Φ to colored particles required for the PQ mechanism. In this subsection, we consider a variant of the Kim-Shifman-Vainshtein-Zakharov (KSVZ) axion model [34,35] in which heavy quarks are charged under $U(1)_{\rm PQ}$ while the SM quarks are not. Matter contents and charge assignments are summarized in Table I. Here, we introduce a doublet scalar field H charged under $SU(2)_H$, in addition to the bifundamental quarks Q, $\bar{Q}$ charged under $SU(3)_c \times SU(2)_H$, for the reason that will become clear below.

TABLE I. Charge assignments in the KSVZ-type model.

	Q	$\bar{Q}$	Φ	H
$SU(3)_c$	$\mathbf{3}$	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\mathbf{1}$
$SU(2)_H$	$\mathbf{2}$	$\bar{\mathbf{2}}$	$\mathbf{3}$	$\mathbf{2}$
$U(1)_Y$	$-2/3$	$2/3$	0	0
$U(1)_{\rm PQ}$	0	-1	1	0

The renormalizable interactions for Φ, $\bar{Q}$, and Q that are allowed by the symmetries are

$$\mathcal{L}_{\rm int} = y {\rm Tr} \bar{Q} \Phi Q + {\rm H.c.} \quad (8)$$

For couplings of order unity, the masses of Q, $\bar{Q}$, $W_H^\pm$, $\Phi^\pm$, and ϕ are of order the SSB scale v. The axion resides in the phase component of Φ_3 for the vacuum (2). Then, Q and $\bar{Q}$ play the role of the PQ quarks that connect the axion to gluons and hidden photons. In the present model, the domain-wall numbers are given by $N_{\rm DW}^{(c)} = 2$ and $N_{\rm DW}^{(H)} = 3$.

We assume that the mass of H, m_H, is parametrically smaller than the SSB scale v. The primary reason to introduce such light doublet $H = (H^{+1/2}, H^{-1/2})$ is to deplete the abundance of W_H and Φ, which would easily exceed the DM abundance by many orders of magnitude. Indeed, the massive $W_H^\pm$ gauge boson and $\Phi^\pm$ scalar boson can decay into the doublet field after the SSB as $W_H^+ \to H^{+1/2} (H^{-1/2})^\dagger$ and $\Phi^+ \to a H^{+1/2} (H^{-1/2})^\dagger$. In the next section, we will estimate the mass of H in order to be consistent with the observed DM abundance.

The PQ quarks are massless before the SSB, and so they are in thermal equilibrium. If the PQ quarks Q and $\bar{Q}$ did not have any interactions with the SM quarks other than the gauge interactions, they would be stable and easily overclose the Universe. To make them unstable, we introduce an interaction between the PQ quarks and the SM quarks as

$$\mathcal{L}_{\rm int,SM} = y_h H Q \bar{d} + {\rm H.c.}, \quad (9)$$

where d collectively represents a right-handed down-type quark in the Standard Model and we have suppressed the flavor index. Note that Q and $\bar{Q}$ are charged under $U(1)_Y$ so that the interaction term is gauge invariant. In the presence of the above interaction, Q and $\bar{Q}$ quickly decay into H and the SM quark after the SSB.

B. DFSZ axion model

Next, we consider a variant of the DFSZ axion model [36,37] in which the SM quarks are charged under the PQ symmetry.

TABLE II. Charge assignments in the DFSZ-type model.

	h_u	h_d	Φ	Ψ	$\bar{\Psi}$
$SU(2)_H$	**1**	**1**	**3**	**2**	$\bar{\mathbf{2}}$
$U(1)_Y$	1	−1	0	0	0
$U(1)_{PQ}$	−1	−1	1	−1	−1

We extend the SM Higgs sector to a two-Higgs-doublet model by introducing h_u and h_d. We also introduce a pair of doublet fermions Ψ and $\bar{\Psi}$ charged under $SU(2)_H$. Matter contents and charge assignments are summarized in Table II.

The interactions for Φ, h_u, h_d, $\bar{\Psi}$, and Ψ that are allowed by the symmetries are

$$\mathcal{L}_{\text{int}} = \lambda_{h2} h_u h_d \text{Tr}\Phi^2 + \frac{y_\Psi}{M_{\text{Pl}}} \text{Tr}\bar{\Psi}\Phi^2\Psi + \text{H.c.} \quad (10)$$

After the PQ symmetry breaking, the hidden fermions Ψ and $\bar{\Psi}$ obtain the mass of order $y_\Psi v^2/M_{\text{Pl}}$, which is around the electroweak scale for $y_\Psi = \mathcal{O}(1)$. We assume that λ_{h2} is sufficiently small so that the interaction term does not affect the successful electroweak symmetry breaking by the Higgs doublets.

The SM quarks play a role of the PQ quarks that connect the axion to gluons. Similarly, integrating out Ψ and $\bar{\Psi}$, the axion is coupled to the hidden photons via the anomaly. In the present model, the domain-wall numbers are therefore given by $N_{\text{DW}}^{(c)} = 6$ and $N_{\text{DW}}^{(H)} = 2$.

The massive $W_H^\pm$ gauge boson and $\Phi^\pm$ scalar boson can decay into the doublet fermion Ψ after the SSB. Note that there is no heavy PQ quark in this model.

III. COSMOLOGICAL DOMAIN-WALL PROBLEM

Axionic cosmic strings are formed when $U(1)_{PQ}$ is spontaneously broken at the SSB. If it were not for the $SU(2)_H$ gauge symmetry and the monopoles, the domain-wall number $N_{\text{DW}}^{(c)} \neq 1$ would imply that there are $N_{\text{DW}}^{(c)}$ degenerate vacua below the QCD scale. After the QCD phase transition, each axionic cosmic string would be attached to $N_{\text{DW}}^{(c)}$ domain walls. Such a string-wall network is stable, and their energy density comes to dominate the Universe at the end of the day, making the Universe intolerably inhomogeneous. However, the domain-wall number and the string-wall dynamics are drastically changed if we correctly take into account the $SU(2)_H$ gauge symmetry and the presence of monopoles in the Universe. In this section, we explain how they solve the cosmological domain-wall problem in the two axion models given in the previous section.

A. Lazarides-Shafi mechanism

Suppose that all degenerate QCD vacua are connected by a local or global symmetry A, which is spontaneously broken at the same time with (or below the energy scale of) the PQ symmetry. In this case, the energetically most favorable configuration of cosmic string is the one around which the PQ phase (i.e., the axion field value divided by f_a) changes by a factor of $2\pi/N_{\text{DW}}^{(c)}$ and the rest of the phase for Φ is complemented by the Nambu-Goldstone (NG) boson of the symmetry A. Such a cosmic string will be attached by only one domain wall due to the QCD instanton effect. In this case, the domain-wall number is effectively reduced to be unity. Then, those topological defects tend to shrink to a point because of the tension of domain walls and are unstable. This is essentially the mechanism proposed by Lazarides and Shafi for the solution to the domain-wall problem [27] (see Ref. [38] for a recent work).

Let us emphasize here that the discrete symmetry that transforms the degenerate vacua into themselves does not have to be embedded into the center of the group A. The necessary and sufficient condition for the mechanism to work is that the degenerate vacua are connected by a symmetry other than $U(1)_{PQ}$. In any case, we call the above solution the Lazarides-Shafi mechanism in this paper.

In the case of the KSVZ-type axion model, the domain-wall number is $N_{\text{DW}}^{(c)} = 2$. The degenerate vacua are transformed into each other via the sign flip: $\Phi \to -\Phi$. This transformation can be realized by the $SU(2)_H$ gauge transformation as well as the $U(1)_{PQ}$ transformation. Hence, the $SU(2)_H$ symmetry plays the role of what we denote as A above, and the degenerated vacua are connected with each other via the (spontaneously broken) $SU(2)_H$ gauge transformation.

As we have seen in the previous section, we have two types cosmic strings: the axionic string and the Alice string. While the axionic string will be attached by two domain walls by the QCD instanton effect, the Alice string will be attached by one domain wall because the $U(1)_{PQ}$ phase rotation around the string is only π [see Eq. (3)]. In particular, the usual axionic string is unstable and separates into a pair of the Alice strings. The cosmological domain-wall problem is therefore avoided by the Lazarides-Shafi mechanism.

In the DFSZ-type axion model, on the other hand, the domain-wall number is $N_{\text{DW}}^{(c)} = 6$, and there are still three physically distinct vacua after taking account of the $SU(2)_H$ gauge symmetry. In this case, one needs another mechanism to avoid the cosmological domain-wall problem, which will be discussed next.

B. Witten effect

Here, we focus on the DFSZ-type axion model, and the cosmological domain-wall problem is intrinsic to this model. Below, we will see the domain-wall problem can be solved by the Witten effect.

As shown in the next section, monopoles form at the SSB of the $SU(2)_H$ symmetry. The axion acquires an effective mass due to the Witten effect in the presence of monopoles even before the QCD phase transition [22–24]. In particular, the $\mathrm{U}(1)_{\mathrm{PQ}}$ symmetry is explicitly broken by the Witten effect down to Z_2 discrete symmetry because of $N_{\mathrm{DW}}^{(H)} = 2$. These degenerate vacua are connected by the $\mathrm{SU}(2)_H$ symmetry, and therefore the effective domain-wall number due to the Witten effect is reduced to unity in combination with the Lazarides-Shafi mechanism. As a result, each cosmic string (of the Alice type) is attached by one domain wall. The tension of those domain walls is determined by the mass of axion as

$$\sigma_a(t) \sim c_a m_{a,M}(t) f_a^2, \tag{11}$$

where $c_a = \mathcal{O}(1)$ is a numerical constant and $m_{a,M}(t)$ is the axion mass due to the Witten effect.[4]

The effective axion mass $m_{a,M}(t)$ is determined by the number density of monopoles as [26]

$$m_{a,M}^2 \simeq 2\kappa \frac{n_M(T)}{f_a}, \tag{12}$$

where

$$\kappa \equiv \left(\frac{N_{\mathrm{DW}}^{(H)}}{N_{\mathrm{DW}}^{(c)}}\right)^2 \frac{\alpha_H}{32\pi^2 r_c f_a}. \tag{13}$$

Here, $r_c = m_\Psi^{-1}$ is the electric screening scale due to Ψ, and κ is of order 10^{-10} for $f_a \simeq 10^{10}$ GeV, $m_\Psi \simeq 1$ TeV, and $\alpha_H \simeq 0.1$. Note that antimonopoles give the same effect with the same sign, so we should regard n_M as the total number density of monopoles and antimonopoles. Domain walls form when the axion mass becomes larger than the Hubble parameter. This happens at the temperature of

$$T_{\mathrm{DW}} \simeq 0.02\ \mathrm{GeV} \left(\frac{\kappa}{10^{-10}}\right)\left(\frac{m_M}{10^{10}\ \mathrm{GeV}}\right)^{-1} \times \left(\frac{\Omega_M h^2}{0.12}\right)\left(\frac{f_a}{10^{10}\ \mathrm{GeV}}\right)^{-1}. \tag{14}$$

This is around the QCD scale, so this effective mass due to the Witten effect can be neglected until the QCD phase transition occurs.

After the QCD phase transition, the axion also acquires a mass from the nonperturbative QCD effects. At zero temperature, it is given by

$$m_a \simeq \frac{\sqrt{m_u m_d}}{m_u + m_d} \frac{m_\pi f_\pi}{f_a}, \tag{15}$$

where m_u, m_d, and m_π are, respectively, the up quark, down quark, and pion masses and f_π is the pion decay constant. The domain-wall number of the axion coupling to gluons is given by $N_{\mathrm{DW}}^{(c)} = 6$, so this nonperturbative effect breaks $\mathrm{U}(1)_{\mathrm{PQ}}$ down to Z_6 discrete symmetry. The degenerate six vacua get into three pairs by the $\mathrm{SU}(2)_H$ symmetry, and the effective domain-wall number is reduced to three by the Lazarides-Shafi mechanism.

Since the two effective domain-wall numbers of the axion couplings to $\mathrm{U}(1)_H$ and $\mathrm{SU}(3)_c$ are relatively prime [$N_{\mathrm{DW}}^{(H)}/2 = 1$ and $N_{\mathrm{DW}}^{(c)}/2 = 3$], the PQ symmetry is completely broken by the Witten effect and the nonperturbative QCD effect. Thus, the vacuum degeneracy is lifted in the presence of monopoles at a scale below the QCD scale. As a result, domain walls experience a pressure from the vacuum energy in the false vacua and shrink to a point, and the entire space will be filled with the true vacuum (which approaches the strong CP-conserving minimum as the Witten effect becomes inefficient). They disappear when

$$\sigma_a(t) = \frac{\mu(t)}{t}, \tag{16}$$

or $T \sim T_{\mathrm{DW}}$, where $\mu(t) \simeq \pi f_a^2 \ln(\delta_s^{-1} t)$ is the tension of the cosmic string and δ_s is its core width. Since T_{DW} is about the QCD scale, the domain wall and cosmic string system disappear soon after the QCD phase transition.

The PQ symmetry breaking effect due to the Witten effect is proportional to the number density of monopoles. Since it decreases inversely proportionally to the volume, its effect is extremely small at the present epoch. It is easy to check that the effective strong CP phase is much smaller than 10^{-10} in our scenario.

IV. DARK MATTER CANDIDATES

In the setup described in Sec. II, the monopole is the only stable object of which the mass is of order or higher than the SSB scale. Below the SSB scale, we have the following stable particles: $\mathrm{U}(1)_H$ charged fields ($H^{\pm 1/2}$ in the KSVZ-type model and $\Psi^{\pm 1/2}$ in the DFSZ-type model), the QCD axion a, and $\mathrm{U}(1)_H$ gauge field in order of decreasing mass. While the $\mathrm{U}(1)_H$ gauge boson contributes to dark radiation, the other particles are massive and contribute to DM. In the following, we estimate the contributions to the DM abundance.

A. Hidden monopoles

We consider the case in which the phase transition is of the first order. This is actually the case in which the minimum of the potential is much larger than the mass scale of the field. This can be realized when λ_1 and λ_2 are sufficiently small. Then, the field is shortly trapped at the origin before the phase transition. In this case, the phase transition ends by the bubble coalescence. The monopole

[4] The effective axion mass depends on the spatial distribution of the monopoles around the cosmic string. Also, the axion potential has a spinodal point at the maximum, and we make only an order of magnitude estimate here.

abundance is determined by the Kibble mechanism [39]: the orientation of the scalar field is random at scales beyond the correlation length, and so monopoles are created with a probability close to unity when the bubbles collide.

The bubble nucleation rate per unit volume per unit time is given by $\Gamma_{\rm bubble} \sim T^4 e^{-S_3/T}$ at a finite temperature T, where S_3 is the three-dimensional Euclidean action of a bubble solution. The phase transition starts when the bubble nucleation rate in a Hubble volume becomes comparable to the Hubble expansion rate $H(t)$,

$$H^4(t_c) \sim T_c^4 e^{-S_3/T_c}. \tag{17}$$

where t_c and $T_c = T(t_c)$ are the time and temperature at the beginning of the bubble nucleation, respectively.

After the nucleation, the bubble wall expands at a velocity close to the speed of light. The phase transition is completed when the entire space is filled with the bubbles. This timescale is given by the inverse of the time derivative of the nucleation rate,

$$\frac{d}{dt} \ln (T_c^4 e^{-S_3/T_c})_{t=t_c} \sim H_c \beta, \tag{18}$$

where we define

$$\beta \equiv -\frac{1}{H} \frac{d}{dt} \frac{S_3}{T}\bigg|_{t=t_c}, \tag{19}$$

and we assume $\beta \gg 1$ in the second equality. In fact, from Eq. (17), we expect

$$\beta \sim \frac{S_3}{T}(t_c) \simeq \ln\left(\frac{T_c^4}{H_c^4}\right), \tag{20}$$

where $H_c \sim T_c^2/M_{\rm Pl}$, and so, for $T_c \ll M_{\rm Pl}$, we find $\beta \gg 1$.

Since the phase transition is completed within the timescale of order $(H_c\beta)^{-1}$ after $t = t_c$, the typical size of the bubbles at the end of phase transition is about $(H_c\beta)^{-1}$. Hence, the number of nucleated bubbles within a Hubble volume is given by $\sim\beta^3$. Since monopoles are created when the bubbles collide, the monopole number density after the phase transition is given by

$$\frac{n_M}{s} \sim \frac{\beta^3 H_c^3}{T_c^3}. \tag{21}$$

Thus, the monopole abundance is given by

$$\Omega_M h^2 \sim 0.1 \left(\frac{m_M}{10^{10}\ \mathrm{GeV}}\right) \left(\frac{T_c}{3\times 10^9\ \mathrm{GeV}}\right)^3 \left(\frac{\beta}{80}\right)^3, \tag{22}$$

where m_M is the monopole mass.

Here, we briefly check if the annihilation of the monopole and antimonopole is negligible. The annihilation cross section is roughly given by $\sigma v_{\rm th} \sim 1/v^2$, where $v_{\rm th}$ is a typical velocity of the monopole. Therefore, the ratio between the annihilation rate and the Hubble expansion rate is roughly given by $\sigma v_{\rm th} n_M/H \sim \beta^3 (T_c^2/v/M_{\rm pl})^4$. This is much smaller than unity for the parameters we are interested in. Therefore, the monopole and antimonopole abundances are determined by Eq. (22). One may wonder if the monopole and antimonopole may form a bound state in the electric plasma because of the attractive Coulomb force and the annihilation cross section may be enhanced. However, this effect is negligible when the phase transition is first order as we consider in this paper [3,40].

To solve the domain-wall problem in the DFSZ-type axion model, the monopole abundance should be sufficiently large [see Eq. (14)]. This puts a lower bound on the scale of the symmetry breaking. The domain walls should disappear before the big bang nucleosynthesis starts. Hence, we require $T_{\rm DW} \gtrsim \mathcal{O}(1)$ MeV, which gives $v(\propto T_c) \gtrsim 10^9$ GeV.

One may consider the Coleman-Weinberg–type potential to realize the strong first-order phase transition. In this case, the breaking scale v is larger than the critical temperature T_c by a factor about $1/g_H$. Hence, we expect that the breaking scale is of order 10^{10} GeV, which is also reasonable for the PQ breaking scale (cf. the next subsection).

B. Nonrelativistic axions

As we have seen in Sec. III, the domain walls and cosmic strings disappear at the QCD phase transition in our models. During the annihilation process of domain walls and cosmic strings, their energy will be released as marginally relativistic axions. Those axions will be nonrelativistic soon after the emission. The resulting energy density of axions can be estimated by [20]

$$\Omega_a h^2 \sim 10^{-2} \left(\frac{f_a}{10^{10}\ \mathrm{GeV}}\right)^{1.19}. \tag{23}$$

This is consistent with the experimental lower bound on the axion decay constant [14], and the PQ breaking scale happens to be close to the monopole mass to explain DM.

C. Charged hidden field

The abundance of $\mathrm{U}(1)_H$ charged fields ($H^{\pm 1/2}$ in the KSVZ-type model and $\Psi^{\pm 1/2}$ in the DFSZ-type model) is determined by the freeze-out mechanism. Here, we focus on the abundance of $H^{\pm 1/2}$, and the result is similar for Ψ (except for the argument about the Higgs portal coupling).

The abundance of H is given by

$$\Omega_H h^2 \approx \frac{5.0\times 10^{-27}\ \mathrm{cm}^3\,\mathrm{s}^{-1}}{\langle \sigma_H v\rangle}, \tag{24}$$

where the annihilation cross section is given by

$$\langle \sigma_H v \rangle \simeq \frac{\pi \alpha_H^2}{m_H^2}. \tag{25}$$

Hence, we obtain

$$\Omega_H h^2 \approx 0.01 \left(\frac{m_H}{1\ \mathrm{TeV}}\right)^2 \left(\frac{\alpha_H}{0.1}\right)^{-2}. \tag{26}$$

In the KSVZ-type axion model, it is possible to include a Higgs portal coupling, $\lambda_H |h|^2 |H|^2$, where h denotes the SM Higgs field. The direct DM search experiments put a very tight constraint on this coupling and exclude most of the parameter space [41] unless the annihilation via the $\mathrm{U}(1)_H$ gauge interaction is efficient. Hence, we assume that the Higgs portal coupling is so small that it does not contribute to the annihilation process.

V. OBSERVATIONAL IMPLICATIONS

A. Kinetic mixing

If the PQ symmetry is classically exact, the kinetic mixing between the $\mathrm{U}(1)_Y$ gauge boson and $\mathrm{U}(1)_H$ gauge boson is absent because of charge conjugation symmetry in the $\mathrm{U}(1)_H$ sector. In the language of high-energy theory, the kinetic mixing will be absent because the operator

$$\frac{\kappa}{M_{\mathrm{Pl}}} \mathrm{Tr} \Phi F_Y F_H + \mathrm{H.c.}, \tag{27}$$

which would give a kinetic mixing of order v/M_{Pl} after the SSB, is prohibited by the PQ symmetry. The above kinetic mixing operator explicitly breaks $\mathrm{U}(1)_{\mathrm{PQ}}$ symmetry, and it contributes to the axion mass as

$$m_{a,\kappa}^2 \sim \frac{\kappa^2}{16\pi^2} M_{\mathrm{Pl}}^2. \tag{28}$$

To solve the strong CP problem, this new contribution must be smaller than about $10^{-10} m_a^2$. Thus, the effective coupling κ in Eq. (27) is severely constrained as $\kappa \lesssim \mathcal{O}(10^{-34})$. One of the possible ways to suppress these dangerous contributions is gauging the discrete Z_N subgroup of $\mathrm{U}(1)_{\mathrm{PQ}}$ [42–45]. In this case, similar explicit PQ breaking operators are prohibited up to $\mathrm{Tr}\Phi^N F_Y F_H$.

If the PQ symmetry is broken by higher-dimensional operators, on the other hand, there appears a small kinetic mixing between $\mathrm{U}(1)_H$ and $\mathrm{U}(1)_Y$. In the rest of this subsection, we focus on the KSVZ-type axion model. Suppose that we have the higher-dimensional operators

$$\frac{\lambda}{M_{\mathrm{Pl}}^n} \mathrm{Tr} \bar{Q} \Phi^{n+1} Q + \mathrm{H.c.}, \tag{29}$$

where n is an integer and λ is a coupling constant of order unity. This term breaks the PQ symmetry and gives an explicit mass to the axion by quantum corrections. The axion mass induced by this operator can be estimated as

$$m_{a,\lambda}^2 \sim \frac{\lambda}{16\pi^2} \left(\frac{v}{M_{\mathrm{Pl}}}\right)^{n-2} v^2. \tag{30}$$

Again, this must be smaller than about $10^{-10} m_a^2$. Noting that $v \sim 10^{10}$ GeV, we find that n should be larger than or equal to 8 to explain the smallness of the strong CP phase.

The same operator violates the mass degeneracy between Q^1 and Q^2: $\delta m \equiv m_{Q^1} - m_{Q^2} = \lambda v (v/M_{\mathrm{Pl}})^{n+1}$. In this case, the one-loop diagrams associated with these fields lead to a kinetic mixing χ between $\mathrm{U}(1)_Y$ and $\mathrm{U}(1)_H$,

$$\chi \sim \frac{g_H g_Y}{16\pi^2} \ln \left(\frac{\delta m_Q}{m_Q}\right)^2, \tag{31}$$

where $m_Q = yv$. The logarithmic factor gives a factor of 10^{-2} or smaller.

At low temperature, the axion VEV cancels the strong CP phase but does not cancel the CP phase associated with $\mathrm{U}(1)_H$ in the hidden sector [see Eq. (7)]. Hence, the monopole generically acquires a hidden electric charge of order unity and becomes a dyon in terms of the $\mathrm{U}(1)_H$ gauge symmetry. As a result, the kinetic mixing induces a mini electric charge for the monopole [46]. The constraint of the kinetic mixing due to the absence of minicharged particles is given by

$$\chi g_H \lesssim 10^{-6}, \tag{32}$$

for the monopole DM with mass of 10^{10} GeV [47,48]. This is satisfied when $g_H \lesssim 0.1$ for $n = 3$.

The charged scalar field $H^{\pm 1/2}$ also acquires the electric charge via the kinetic mixing. If its amount is as large as the observed DM abundance, the null result of the minicharged particle search gives a stronger constraint by a factor of 10^{-4}. Since the kinetic mixing depends on the PQ breaking parameter only logarithmically, this constraint excludes the scenario with even a tiny amount of the explicit PQ breaking effect. Therefore, the darkness of DM is closely tied to the quality of PQ symmetry in this case. The high quality of PQ symmetry may be explained by the anthropic argument because the large-scale structure does not form in the Universe in a scenario of charged DM.

In the case of the DFSZ-type axion model, there are no particles that are charged under both $\mathrm{U}(1)_H$ and $\mathrm{U}(1)_Y$. Then, the kinetic mixing comes only from the higher-dimensional operator of Eq. (27).

B. Dark radiation

The axion is thermalized after the PQ phase transition [49,50]. The remaining hidden gauge field $\mathrm{U}(1)_H$ as well as the axion thus contribute to the energy density of the

Universe as dark radiation. The amount of dark radiation $\rho_{\rm DR}$ is conveniently described by the effective neutrino number $N_{\rm eff}$ as

$$\Delta N_{\rm eff} = \frac{4}{7}\frac{\rho_{\rm DR}}{(\pi^2/30)T_\nu^4}, \tag{33}$$

where T_ν is the neutrino temperature.

First, we consider the KSVZ-type axion model. In the absence of the Higgs portal interaction, the hidden gauge field $\mathrm{U}(1)_H$ and $H^{\pm 1/2}$ are thermalized and are decoupled from the SM sector well before the electroweak phase transition. Then, the latter field annihilates into the former one during the freeze-out epoch. The resulting energy density of the hidden gauge field is thus larger than that of the axion by a factor of $2\times 3^{4/3}$. Hence, we obtain

$$\Delta N_{\rm eff} = \frac{4}{7}(2\times 3^{4/3}+1)\left(\frac{g_*(T_D)}{43/4}\right)^{-4/3} \tag{34}$$

$$\simeq 0.26, \tag{35}$$

where g_* $(=106.75)$ is the effective relativistic degrees of freedom in the standard model sector at the decoupling temperature T_D.

If the Higgs-portal interaction is strong enough, the hidden sector is in thermal equilibrium around the electroweak scale until the density of the charged scalar $H^{\pm 1/2}$ freezes out. In this case, the decoupling temperature is about $m_H/10$, and g_* is about 90–100 for $m_H \gtrsim 100$ GeV. Since the annihilation into the $\mathrm{U}(1)_H$ gauge boson dominates, the energy density of hidden gauge field is again larger than that of the axion by a factor of $2\times 3^{4/3}$ except for the difference of g_*. Hence, we obtain

$$\Delta N_{\rm eff} = \frac{4}{7}\left(\frac{106.75}{43/4}\right)^{-4/3} + \frac{4}{7}2\times 3^{4/3}\left(\frac{90-100}{43/4}\right)^{-4/3}$$
$$\simeq 0.28-0.32. \tag{36}$$

The Planck data combined with the observation of baryon acoustic oscillations gives the constraint $N_{\rm eff} = 3.15\pm 0.23$ [51]. The standard model prediction is $N_{\rm eff} = 3.046$, and hence our result is consistent with the constraint. It is expected that the ground-based stage-IV cosmic microwave background (CMB) polarization experiment (CMB-S4) will measure $N_{\rm eff}$ with a precision of $\Delta N_{\rm eff} = 0.0156$ within 1σ level [52] (see also Ref. [53]).

Next, we consider the DFSZ-type axion model, in which the hidden gauge field and $\Psi^{\pm 1/2}$ are decoupled from the SM sector well before the electroweak phase transition. The effective relativistic degrees of freedom of the charged fermions is given by $2\times 4\times 7/8$ at a high temperature. After the freeze-out epoch, the effective number of neutrinos is given by

$$\Delta N_{\rm eff} = \frac{4}{7}(2\times (9/2)^{4/3}+1)\left(\frac{g_*(T_D)}{43/4}\right)^{-4/3} \tag{37}$$

$$\simeq 0.42, \tag{38}$$

where we used $g_* = 106.75$. This is consistent with the Planck constraint within 2σ level.

VI. DISCUSSION AND CONCLUSIONS

Motivated by the coincidence of the energy scales, we have pursued the possibility of unifying the PQ symmetry breaking and the production of the monopole DM. We have provided both KSVZ- and DFSZ-type axion models, in which the cosmological domain-wall problem can be avoided either by a mechanism in the same spirit of the Lazarides-Shafi mechanism or by the Witten effect.

An SU(2) doublet field needs to be introduced to make some unwanted heavy relics unstable. As a result, there are three candidates of DM: the axion, monopole, and hidden-charged field. The latter two are charged under the remaining U(1) gauge symmetry, so they may acquire a nonzero electric charge via a possible kinetic mixing between the electroweak and hidden U(1) gauge bosons. We have found that the amount of kinetic mixing, or the electric charges of that DM, is related to the quality of the PQ symmetry. Hence, the darkness of DM, which is required by the large-scale structure formation, may explain the high quality of the PQ symmetry.

The hidden gauge bosons, as well as the relativistic components of axions, contribute to the energy density of the Universe as dark radiation. The amount of that energy density is consistent with the present observational result and can be distinguished from the standard scenario in the near future.

Finally, we comment on the detectability of gravitational waves that are emitted from the dynamics of the bubble at the first-order phase transition. Since the energy scale of the phase transition is as high as 10^{10} GeV, the typical frequency of those gravitational waves is too high to be detected by the proposed detectors [54].

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